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PLASMA INSTABILITIES AND MODEL RADIO AURORA

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16. Abstract

The totality of theoretical work dealing with the possible types of plasma instabilities associated with currents, fluxes of incoming particles, and electron density gradients is discussed in the present review with reference to the polar ionosphere. The theory developed up to the present time explains in general terms the properties of the observed radio reflections.

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Introduction

Already in the 1950's, Lovell, Clegg, and Ellyett [1947] discovered radio reflections associated with the polar aurorae during radar observations of the E-regions of the polar ionosphere. The essence of this phenomenon lay in the fact that, simultaneously with magnetic disturbances, there arise small scale inhomogeneities in the electron level extended along the Earth's magnetic field in which ultrashort waves of the 30 -- 1000 MHz region are strongly scattered. Such a specific phenomenon was immediately tabbed the "radio aurora." The first systematic investigations of radio aurora were begun in the 1960's [Moore, 1951; McNamara, 1955; Bullought, 1954; Collins, 1951]. /3*

The distinguishing peculiarity of these radio reflections is the sharp dependence of the scattered signal intensity on the angle between the normal to the front of the incident wave and the plane perpendicular to the Earth's magnetic field vector in the scattering region. This angle has been called the aspect angle ψ. The radio reflections possess the fundamental property of aspect sensitivities. This property consists of the fact that the reflections occur principally from those regions where the directional vector $\vec{R}$ of the radiowave is orthogonal to the force line of the geomagnetic field $\vec{B}$, i.e., when

*Numbers in the margin indicate the pagination of the original foreign text.

$$\vec{R}\vec{B} \approx 0 \qquad (1)$$

or the aspect angle ψ is close to zero.

The general geometry of the radio reflections is schematically illustrated in Figure 1. According to the communications of Leadabrand [1965], the reflected signal intensity sharply decreases by more than 10 dB at a deviation by an angle of the order of 1° from the orthogonal direction; however, in a series of other cases, the condition of orthogonality is not so rigorous [Hofstee, Forsyth, 1969; Moorcroft, 1972a; Pyatsi, Sverdlov, 1971].

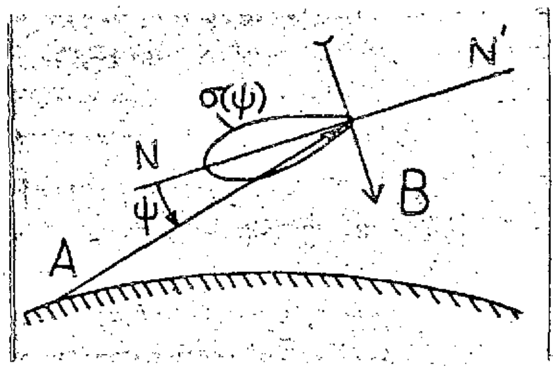

Figure 1. Schematic illustration of the auroral radio reflections.

A- location of the transmitter; B- vector of the Earth's magnetic field, N, N'- line orthogonal to the magnetic field; ψ- aspect angle; $\sigma(\psi)$- effective scattering cross section as a function of the aspect angle.

Various suggestions as to the mechanism of formation of the inhomogeneities and their scattering of radiowaves were advanced by Moore (1952), Chapman (1953), Booker (1956, and Herlofson (1947). The assumption that the reflecting centers of "auroral ionization" are distributed along the force line of the Earth's magnetic field, forming extended structures along them, lies at the basis of the work of Herlofson and Chapman. The geometrical locations satisfying Equation (1) are computed by Chapman, and thus the configurations of the reflecting regions were derived for various observation points. Although Chapman's geometrical scheme appears to be an effective tool, it contains, however, no statement as to the origin of the anisotropy of the ionization. Chapman and Herlofson considered it to be self-evident that the ionization along the magnetic field force lines can immediately be produced by the motion of incoming particles.

The first attempt to explain the anisotropy of the ionization was proposed by Booker (1956). According to Booker's proposal, the radio reflections are caused by noncoherent scattering by weak ($\Delta N/N \sim 10^{-3} - 10^{-4}$) turbulent inhomogeneities of the electron level which are shaped like columns extended along the magnetic field which have long dimensions of $L \sim 10 - 20$ m and transverse dimensions of $T \sim 3 - 5$ m and are distributed at random in the scattering volume.

In Booker's model, the ionosphere is treated as a liquid /6 in a turbulent condition; therefore, the physical premises which lie at its basis are already at odds with contemporary ideas. Later, with the accumulation of experimental data, Booker's model was improved and generalized [Egeland, 1963; Moorcroft, 1961]. Thus, Egeland generalized Booker's theory, introducing ellipsoidal inhomogeneities into the discussion, and Moorcroft introduced an ensemble of Gaussioids of various sizes and critical frequencies. This improvement permitted the introduction of the mechanism of critical reflection for the lower frequencies and partial reflection for the higher frequencies with a transition from one region to the other.

However, later the accumulated extensive volume of experimental data revealed their incompatibility with the cited improved models.

For the first time, the idea of plasma turbulence, as this term is now understood, was advanced by Farley (1963) and Buneman (1963) to explain the radio reflections in the equatorial E-layer of the ionosphere. It was asserted in these papers that

the small scale inhomogeneities which reflect the radar signal of a superposition of the plane density waves of charged particles, arise as a result of the development of an instability in the ionospheric uniform plasma, with the current of the relatively low frequency longitudinal disturbances. Later, ideas about plasma turbulence began to attract attention as an explanation of the origination of radio reflections in the regions of polar aurora [Leadabrand and Hodges, 1965; McDiarmid and McNamara, 1969; Hofstee and Forsyth, 1969; Unwin and Knox, 1971; Moorcroft, 1972]. In later papers of Unwin and Baggaley (1972), Knox (1972), /7
Ogowa, Sato (1972), and Sato (1971), the possibility was discussed of explaining auroral radio reflections as turbulent pulsations arising as a result of the development of a drift gradient instability.

At the present time, the term 'turbulence' is of a more general nature than twenty years ago and refers to liquids, plasmas, and solid objects. According to contemporary notions, the usual turbulent state in a plasma is the result of the development of various types of instabilities.

A weakly ionized plasma in a magnetic field with gradients in the ionized density present along with electric fields becomes the source of many types of instabilities which have not yet been completely investigated. The various types of instabilities theoretically well studied in a laboratory plasma [Buneman, 1963; Simon, 1963; Hoh, 1963; Kadomtsev, 1964; D'Angelo, 1965] have been considered with regard to the ionosphere in the papers of [Farley, 1963; Reid, 1968; Tsuda and Sato, 1966; Knox, 1964; Rogister and D'Angelo, 1970; Perkins, 1968] and appeared to be the source of small scale turbulence.

If, under the influence of external effects (electric field, currents, and so forth), the waves in a certain frequency and wave number interval become unstable in a plasma, and their energy reaches a level sufficient for the appearance of nonlinear processes transforming energy from the production region to the absorption region, then the establishment of steady state turbulence is possible. A balance between production and absorption occurs. If W_K^σ is the spectral energy density of waves of the σ type, the balance equation is of the general form:

/8

$$\frac{\partial W_{\vec{K}}^\sigma}{\partial t} + \vec{V}_g \frac{\partial W_{\vec{K}}^\sigma}{\partial \vec{r}} = \gamma(J_{\vec{K}}^\sigma) W_{\vec{K}}^\sigma + \int \Phi(W_{\vec{K}}^\sigma, W_{\vec{K}'}^\sigma)\, d\vec{K}' \quad (2)$$

The first term on the left is determined by the explicit time dependence of the spectral densities, and the second term reflects the energy transfer of the plasma waves in space. The first term on the right describes the linear production of σ-type waves and the possible quasilinear limitation on the production, and the second term on the right describes the nonlinear transfer of turbulence energy over the spectrum in the case of processes of the induced scattering type, discharges, and so on. Upon the calculation of the nonlinear effects of just the second order in the turbulent energy, the function $\Phi(W_{\vec{K}}^\sigma, W_{\vec{K}'}^\sigma)$ is quadratic in the second term with respect to the turbulent energy.

In the case $\frac{\partial W_{\vec{K}}^\sigma}{\partial t} = 0$, the dependence of $W^\sigma(\vec{K})$ determines the quasistationary spectrum of the turbulence, which in the final accounting is responsible for the radio reflections. Therefore, the clarification of the possibility of different types of instabilities in the polar ionosphere and the discovery of the quasistationary turbulent spectra to which they lead prove to be important. Below, we dwell on a discussion of the linear theory of instability, which makes it possible to determine the conditions under which turbulence arises, the buildup increments, and the frequencies $\omega(\vec{K})$ of the unstable waves. The linear theory does not, however, answer the question of whether or not the mechanism discussed can maintain the observed radio reflection intensities, and it does not give the spectrum $W_{\vec{K}}^{\epsilon}$ of the turbulence considered, which is responsible for the radio reflections. Only the nonlinear theory, which is being intensively developed at the present time, can provide an answer to the last question.

§ 1. Buneman-Farley Instability

As has already been pointed out, Buneman (1963) and Farley (1963) proposed representing small scale inhomogeneities oriented along the Earth's magnetic field as a superposition of plane density waves of charged particles arising as a result of an instability in the ionospheric plasma with the current of relatively low frequency longitudinal disturbances. /9

The large amounts of experimental data regarding magnetic disturbances in the polar ionosphere show that strong currents corresponding to a drift velocity of the electrons of up to 2 km/sec can occur in the region of the E ionosphere, according to the data of [Boestrom, 1964; Mozer and Bruston, 1967; and Wentworch and Potter, 1970]. If the drift velocity of the electrons relative to the ions exceeds a certain threshold value of the order of the ion thermal velocity, then specific iono-sonic waves will be produced in such a system. In order to avoid confusions with the usual iono-sonic wave in a nonisothermal plasma, these waves will be referred to as Farley waves in what follows. Waves with wave vectors $\vec{K}$ perpendicular to the magnetic field $\vec{B}$ will intensify, and waves in other directions will be strongly suppressed due to Landau damping.

The Buneman-Farley instability develops in a weakly ionized plasma due to the difference for electrons and ions of the ratio of the gyrofrequencies to the frequency of collision with neutral particles, i.e., $\omega_{He}/\nu_e \gg 1$ and $\omega_{Hi}/\nu_i \ll 1$. The drift of electrons is proportional to $V_d = \frac{c[EB]}{B^2}$, and in the case of practically non-drifting ions, it gives a two-stream instability. Buneman has /10 shown that such an instability can excite disturbances oriented along the field and propagated in the direction $\vec{E} \times \vec{B}$ in the ionosphere [Buneman, 1963]. Buneman's idea was further developed by Farley, who applied the Boltzmann-Vlasov kinetic equation in order to derive the kinetic drift velocity and the conditions for an instability to arise.

In order to clearly show the physical meaning of the instability being discussed, let us introduce the example of the discovery of the dependences $\omega(\kappa)$ and $\gamma(\kappa)$ for Farley-waves in a uniform plasma, using a simplified hydrodynamic model. We will proceed from the linear hydrodynamic equations of motion of charged particles for electrons and ions, the continuity equations, and the Poisson equations:

$$\begin{aligned}
&m\frac{\partial \vec{V}_e}{\partial t} = -e\left(\vec{E} + \frac{1}{c}[\vec{V}_e \cdot \vec{B}]\right) - m\nu_e\vec{V}_e - \frac{V_{Te}^2\,\vec{\nabla} n_e}{n_e}m,\\
&M\frac{\partial \vec{V}_i}{\partial t} = e\left(\vec{E} - \frac{1}{c}[\vec{V}_i \cdot \vec{B}]\right) - M\nu_i\vec{V}_i - \frac{V_{Ti}^2\,\vec{\nabla} n_i}{n_i}M,\\
&\frac{\partial n_e}{\partial t} + \vec{\nabla}(n_e\vec{V}_e) = 0, \qquad \frac{\partial n_i}{\partial t} + \vec{\nabla}(n_i\vec{V}_i) = 0,\\
&n_{e0} = n_{i0}, \quad V_{Te}^2 = \frac{T_e}{m}, \quad V_{Ti}^2 = \frac{T_i}{M}, \quad \operatorname{div}\vec{E} = 4\pi e(n_i - n_e).
\end{aligned} \tag{3}$$

In the polar ionosphere, the conditions mentioned $\nu_e \ll \omega_{He}$ and $\nu_i \gg \omega_{Hi}$ are fulfilled in the altitude interval 120 > h > 80 km with the use of a model ionosphere [Alpert, 1972]. Since we are seeking unstable low frequency oscillations, we can immediately neglect the inertia of the electrons. Let us also neglect the term $\frac{1}{c}[\vec{V}_i \cdot \vec{B}]$ in the equation of motion of the ions, i.e., the drift of the ions associated with the magnetic field.

Then the equation of motion for the electrons can be written in the form

$$\vec{V} = \frac{e}{m\nu_e}\left(\vec{E} + \frac{1}{c}[\vec{V}_e \cdot \vec{B}]\right) - \frac{V_{Te}^2\,\vec{\nabla} n_e}{n_e \nu_e}, \tag{4}$$

and for the unperturbed steady state condition: /11

$$\vec{E}_0 + \frac{1}{c}[\vec{V}_0 \cdot \vec{B}] = \frac{m\nu_e}{e}\vec{V}_0 \tag{4a}$$

With the condition that the fields $\vec{E}$ and $\vec{B}$ cross, the solution of Equation (4a) can be obtained in the form:

$$\vec{V}_{e0} = \frac{c}{B_0^2}[\vec{E}_0 \cdot \vec{B}] - \frac{\nu_e c \vec{E}_0}{\omega_{He} B}, \tag{5}$$

where the last term of Equation (5), which is proportional to the collisional frequency, is small due to the condition $\nu_e \ll \omega_{He}$ and determines the motion along the electric field. Thus, the electrons essentially drift orthogonally to the electric and magnetic fields, and due to collisions, they are somewhat smeared out in the directions opposite to the electric field. The ions in the steady state condition with $\nu_i \gg \omega_{Hi}$ have a drift velocity of

$$V_{0i} = \frac{eE_0}{\nu_i M} = \frac{\omega_{Hi}}{\nu_i} c \frac{E_0}{B_0} \ll V_{0e},$$

the consideration of which in the discussion of the oscillations gives nonessential corrections. The coordinate system is illustrated in Figure 2.

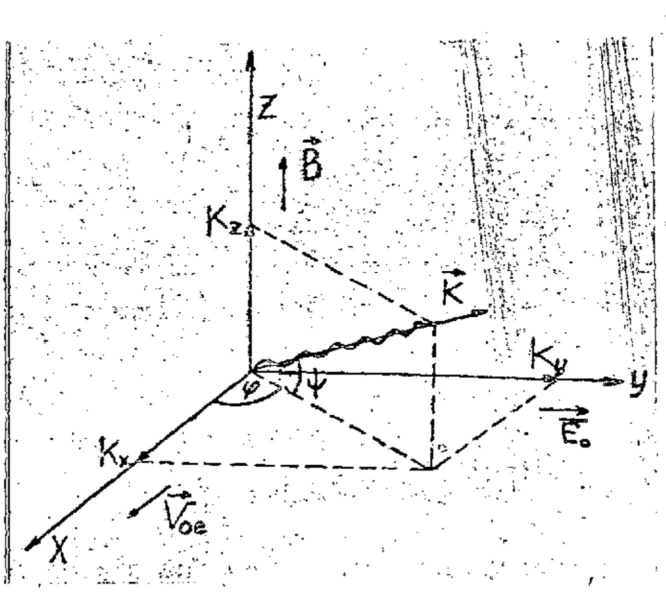

Figure 2. Coordinate system.

Furthermore, let us introduce the linearization procedure setting

$$\begin{vmatrix} \vec{E} = \vec{E}_0 + \vec{E}_1 , \\ \vec{V} = \vec{V}_0 + \vec{V}_1 , \\ n = n_0 + n_1 , \end{vmatrix}$$

where $\begin{vmatrix} \vec{E}_1 , & \vec{V}_1 , & n_1 \end{vmatrix}$ are quantities of the first order of smallness and are porportional to

$$\begin{vmatrix} exp\,(i\omega t - i\vec{k}\vec{r}) \end{vmatrix}$$

From Equation (3), we obtain for each harmonic the system of equations: /13

$$\begin{vmatrix} -\frac{e}{m}\left(\vec{E}_1 + \frac{1}{c}[\vec{V}_{e1}\vec{B}]\right) - \vec{V}_{e1}\nu_e + \frac{i\vec{K}V_{Te}^2 n_{e1}}{n_0} , \\ i\omega\vec{V}_i = \frac{e}{M}\vec{E}_1 - \nu_i\vec{V}_{i1} - \frac{i\vec{K}V_{Ti}^2 n_{i1}}{n_0} , \\ i\omega n_{e1} - i n_0\vec{K}\vec{V}_{e1} - i n_{e1}\vec{K}\vec{V}_{0e} = 0 , \\ i\omega n_{i1} - i n_0\vec{K}\vec{V}_{i1} - i n_{i1}\vec{K}\vec{V}_{0i} = 0 . \end{vmatrix} \tag{6}$$

Solving the last two equations gives

$$\begin{vmatrix} n_{e1} = n_0\frac{\vec{K}\vec{V}_{e1}}{\omega - \vec{K}\vec{V}_0} , \quad n_{i1} = \frac{n_0\,\vec{K}\vec{V}_{i1}}{\omega - \vec{K}\vec{V}_{0i}} \simeq \frac{n_0\,\vec{K}\vec{V}_{i1}}{\omega} , \end{vmatrix}$$

and we obtain

$$0 = -\frac{e}{m}\left(\vec{E}_1 + \frac{1}{c}[\vec{V}_{e1}\cdot\vec{B}]\right) - \nu_e\vec{V}_{e1} + \frac{i\vec{K}V_{Te}^2(\vec{K}\vec{V}_{e1})}{\omega - \vec{K}\vec{V}_0} , \tag{7}$$

$$0 = \frac{e}{M}\vec{E}_1 - (\nu_i + i\omega)\vec{V}_{i1} - \frac{i\vec{K}V_{Ti}^2(\vec{K}\vec{V}_{i1})}{\omega} \tag{8}$$

Upon the assumption that the waves are longitudinal, i.e., $\vec{K} \parallel \vec{E}_1$, and for the particular case of waves propagating strictly orthogonally to the magnetic field ($\psi = 0$), and, accordingly, $E_{1z}=0$, we obtain, having set $\vec{K} = \vec{i}k$ and $\vec{E} = \vec{i}E$, the equation for the components of the electron velocity $V_{ex} V_{ey}$:

$$0 = -\frac{e}{m}E_1 - \omega_{He} V_{ey} - \nu_e V_{ex} + ikV_{Te}^2 \frac{kV_{ex}}{\omega - kV_{0}} \tag{9}$$

$$0 = \omega_{He} V_{ex} - \nu_e V_{ey} \tag{10}$$

It is evident from Equation (10) that the velocity component V_{ex} along the electric field is not equal to zero only because of the collisions of electrons with neutral particles, $\nu_e \neq 0$. It is appropriate to point out here that the collisions of /14
electrons with neutral particles are necessary for the existence of Farley waves and do not only give the usual dissipation corrections.

Substituting from Equation (9) $V_{ey} = V_{ex}\frac{\omega_{He}}{\nu_e}$ into Equation (8), we obtain

$$0 = -\frac{e}{m}E - \frac{\omega_{He}^2}{\nu_e} V_{ex} - \nu_e V_{ex} + ik^2V_{Te}^2 \frac{V_{ex}}{\omega - kV_{e0}} \tag{11}$$

from which

$$V_{ex} = \frac{-i\frac{e}{m}E}{\frac{\omega_{He}^2}{i\nu_e} - \frac{k^2V_{Te}^2}{\omega - kV_{e0}}} \tag{12}$$

We emphasize that in contrast to the usual ionic sound, in which the disturbed motion of the electrons occurs only along the electric field, both the parallel E_1 velocity component and the perpendicular V_{ey} velocity component, with $V_{ey}/V_{ex} = \frac{\omega_{He}}{\nu_e} \gg 1$ are present in the iono-sonic Farley waves.

For ions, we obtain from Equation (8):

$$V_{i1x} = \frac{-i\frac{e}{M}E_1}{(\omega - i\nu_i) - \frac{K^2 V_{Ti}^2}{\omega}} \tag{13}$$

Substituting the expressions for n_{e1}, n_{i1} into the Poisson equation $\mathrm{div}\,\vec{E} = 4\pi e(n_i - n_e)$, we obtain:

$$-i(\vec{K}\vec{E}_1) = 4\pi e n_0 \left(\frac{K V_{i1x}}{\omega} - \frac{K V_{e1x}}{\omega - \vec{K}\vec{V}_{0e}}\right). \tag{14}$$

After substitution into the Poisson equation (14) of the expressions for V_{e1x}, V_{i1x} from Equations (12) and (13), we obtain the dispersion equation for low frequency longitudinal waves:

$$1 = \frac{\omega_{0i}^2}{\omega(\omega - i\nu_i) - K^2 V_{Ti}^2} + \frac{\omega_{0e}^2}{\frac{\omega_{He}^2}{i\nu_e}(\omega - \vec{K}\vec{V}_{0e}) - K^2 V_{Te}^2}, \tag{15}$$

where /15

$$\omega_{0i}^2 = \frac{4\pi e^2 n_0}{M}, \quad \omega_{0e}^2 = \frac{4\pi e^2 n_0}{m}.$$

We will solve Equation (15) by the method of successive approximations, having selected as the real zero$^{\text{th}}$ approximation $\omega = \vec{K}\vec{V}_0$. In the next approximation, we will obtain $\omega = \omega' - i\gamma$, where

$$\omega' = \frac{\vec{K}\vec{V}_0}{1 + \frac{\nu_i \nu_e}{\omega_{Hi}\omega_{He}}}, \tag{16}$$

$$\gamma = \frac{\nu_e}{\omega_{Hi}\omega_{He}}\left((KV_0)^2 - K^2\frac{T_e + T_i}{M}\right). \tag{17}$$

Equations (16) and (17) agree to an accuracy of the corrections of the next order with the corresponding values derived in the papers of [Kamenetsaya, 1969; Kato and Hirata, 1967; Waldtefel, 1965; Rogister and D'Angelo, 1970]. One can determine from Equation (17) the critical value of the drift velocity at which instability begins:

$$V_0 > V_{crit} \left| \sim \sqrt{\frac{T_e + T_i}{M}} \right| \qquad (18)$$

Consequently, the phase velocity of Farley waves is given by

$$V_F = \left| \frac{\omega}{K} = \frac{(\vec{K}\vec{V}_0)}{|K|} \frac{1}{1 + \frac{\nu_e \nu_i}{\omega_{He}\omega_{Hi}}} \right|$$

and is less than the drift velocity of the particles.

For the case of waves propagating not strictly perpendicularly to the magnetic field (i.e., the angle ψ between $\vec{K}$ and the plane perpendicular to $\vec{B}$ is not equal to zero), the problem of finding the dispersion law and the increment is seriously complicated mathematically. This difficulty is associated with the fact that upon the propagation of waves not orthogonal to the magnetic field, the waves are not rigorously separated into longitudinal and transverse, and consequently, there occurs in each wave both longitudinal and transverse components. It is simplest of all to derive the dispersion equation in the case of $\psi \neq 0$ in the following manner. It follows from Maxwell's equations that /16

$$\frac{c^2 K^2}{\omega^2}\left(\vec{E}_1 - \frac{\vec{K}(\vec{K}\vec{E}_1)}{K^2}\right) = \vec{E}_1 - i\frac{4\pi}{\omega}\vec{j}, \qquad (19)$$

where

$$\vec{j} = en_0\left(\vec{V}_{i1} n_{i0} - \vec{V}_{e0} n_{e1} - \vec{V}_{e1} n_{0e}\right) \qquad (20)$$

(V_{io} is assumed to be negligibly small). Having substituted into Equation (19) the expressions for $|\vec{V}_{i1}|$ found from the equations of motion of the ions (8), having resolved both equations into components in advance, one can express with the aid of these

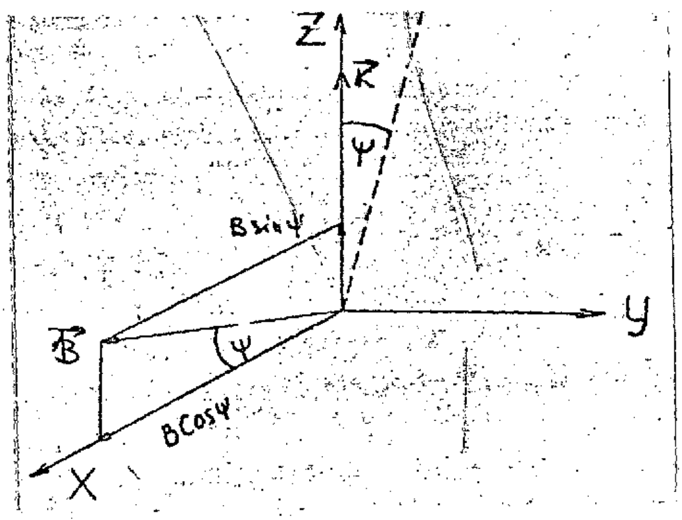

Figure 3. Coordinate system.

two equations, the values of the electric field components E_{ix}, E_{iy}, E_{iz} as a function of the electron velocity components. Thus, it is most convenient to select the coordinate system so that the z-axis is directed along the wave vector $\vec{K}$, and the vector $\vec{B}$ lies in the xz plane (see Figure 3).

Substituting the electric field components E_{ix}, E_{iy}, E_{iz} into the equation of motion of the electrons and equating the determinant of the system of three equations obtained for $\vec{V}_{ie}$ to zero, we find the dispersion equation for Farley waves, taking into account $\psi \neq 0$. The procedure described in the present case is shorter than that generally adopted, in which, on the contrary, the components of the conductivity tensor σ_{ij} (or ε_{ij}) are found from the equations of motion of the electrons and the ions and further, the determinant of the system for the electric field E_i is considered. Corrections derived by taking into account the not strictly longitudinal nature of the waves are always negligibly small, as estimates show, for Farley waves in the ionosphere

E region. The components of the electric field transverse to the propagation direction also do not play a significant role. /18
The application of any of the methods gives a dispersion equation

$$1=\frac{\omega_{0i}^2}{\omega(\omega-i\nu_i)-k^2V_{Ti}^2}+\frac{\omega_{0e}^2\left(1+\frac{\omega_{He}^2}{\nu_e^2}\psi^2\right)}{\omega(\omega-\vec{k}\vec{V}_0)\frac{\omega_{He}^2}{i\nu_e}-k^2V_{Te}^2\left(1+\frac{\omega_{He}^2}{\nu_e^2}\psi^2\right)} \tag{21}$$

Here we have assumed the angle ψ to be small, so that $\sin\psi\sim\psi, \cos\psi\sim 1$. Expressions for the frequency and buildup increment of the wave corresponding in accuracy as far as the correction with the equations derived by Kamenetskaya (1969) are easily derived by the method of successive approximations in the following forms:

$$\omega=\frac{(\vec{k}\vec{V}_0)}{1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}\left(1+\frac{\omega_{He}^2}{\nu_e^2}\psi^2\right)}, \tag{22}$$

$$\gamma=\frac{\nu_e}{\omega_{He}\omega_{Hi}}\left(1+\frac{\omega_{He}^2}{\nu_e^2}\psi^2\right)\left[\frac{(\vec{k}\vec{V}_0)^2}{1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}\left(1+\frac{\omega_{He}^2}{\nu_e^2}\psi^2\right)}-k^2V_{Ti}^2\right] \tag{23}$$

From this follows the criterion for an instability to arise:

$$V_0^2>V_{Ti}^2\left(1+\left(\psi^2+\frac{\nu_e^2}{\omega_{He}^2}\right)\left(\frac{M}{m}\frac{\nu_i}{\nu_e}\right)\right). \tag{24}$$

We note that Equation (24) does not give the dependence of the threshold value $\vec{V}_0$ on $\vec{K}$ because viscosity was neglected. A more detailed analysis of the dispersion laws and the condition for exciting Farley waves can be carried out only on the basis of the kinetic equations, which was first done in Farley's paper (1963). A system of kinetic equations for α-particles has the form

$$\frac{\partial f_\alpha}{\partial t}+\vec{V}_\alpha\vec{\nabla}_r f_\alpha+\frac{e}{m}\left(\vec{E}+\frac{1}{c}[\vec{V}_\alpha\vec{B}]\right)\vec{\nabla}_v f_\alpha=-\nu_{\alpha m}\left(f_\alpha-f_{\alpha m}\int f_\alpha d\vec{V}_\alpha\right),$$

$$f_{\alpha m}=\left(m/2\pi T_{\alpha m}\right)^{3/2}\exp\left(-\frac{m_\alpha\vec{V}_\alpha^2}{2T_{\alpha m}}\right).$$

$$\operatorname{div} \vec{E} = 4\pi N \sum_{\alpha=e,i} e_\alpha \int f_\alpha \, d\vec{v}_\alpha ,$$ /19

$$\operatorname{rot} \vec{E} = -\frac{1}{c} \frac{\partial \vec{B}}{\partial t}$$

Farley solved this system of equations numerically.

The following conclusions, important for the diagnostics of an instability in the ionosphere (see the coordinate system in Figure 2a), can be drawn from consideration of Farley's results:

1. An instability occurs when the drift velocity of the electrons exceeds some threshold value of the order of V_{Ti} 0.1 V_{Te} for waves with a wave vector orthogonal to $\vec{B}$.

2. The critical drift velocity necessary for the excitation of waves rapidly increases as the aspect angle ψ deviates from zero.

3. The width of the angular spectrum increases as the drift velocity grows and is usually of the order of 1 — 2°.

4. The phase velocities of the waves are close to the local sound velocity in the E region of the order of 300 — 500 m/sec.

5. The probability of reflections depends on the azimuthal angle in the form cos ψ.

In addition, one should emphasize especially three more conclusions:

6. The larger the critical drift velocity is, the greater the wave number K.

7. At one and the same drift velocity, the angular spectrum is somewhat narrower for larger wave numbers K.

8. The phase velocities of the waves increase somewhat as K increases.

The main deficiency of Farley's work was the fact that the solution was derived numerically, and therefore it is difficult to investigate in a general way the dependences of the results on such parameters as the collision frequency, wave number, temperature, electrons and ions, and so forth. An analytical solution of the system of kinetic equations simultaneously with Maxwell's equations was derived for the first time in the paper [Kamenetskya, 1970]. The analytic solution with the kinetic corrections is similar to Equations (22) and (23) and has the form: /20

$$\omega = \frac{\vec{K}\vec{V}_0}{1+R}, \quad R = \frac{\nu_i}{\nu_e}\frac{M}{m}\left(\sin^2\psi + \frac{\nu_e^2}{\omega_{He}^2}\right)(1-\eta_i);\ \eta_i = \frac{3K^2T_i}{\nu_i^2 M},$$
$$\gamma = \frac{K^2 R}{\nu_i(1+R)^3}\left[V_{0K}^2\left(1+\frac{\nu_i}{\nu_e}R\right) - \frac{T_i + T_e(1-\eta_i)}{M}(1+R)^2 - 2\eta_i V_{0K}^2\right], \tag{25}$$

where V_{ok} is the projection of the drift velocity $\vec{V}_o$ in the $\vec{K}$ direction.

It is easy to see from the equation for the increment that, upon the condition of a small supercriticality, i.e., when inequality $(V_{0K} - V_s) \ll V_0$, where $V_s^2 = \frac{T_e + T_i}{M}$, the increment drops off sharply to zero upon an increase of ψ due to the multiplier $\left(\sin^2\psi + \frac{\nu_e^2}{\omega_{He}^2}\right)$. In the case of a sufficiently large supercriticality, two small maxima are observed in the dependence of the increment on the angle ψ at $\psi \sim \pm\psi^*$, and a minimum is observed at ψ_o. The instability develops when the drift velocities exceed the critical value:

$$V_{0k}^2 > \frac{T_i + T_e(1-\eta_i)}{M} \frac{(1+R)^2}{(1+\frac{\nu_i}{\nu_e}R)} \tag{26}$$

The last expression refines Equation (17).

We note the kinetic corrections to the solutions are significant only for high frequencies, at which they provide a decrease of the increment. The conclusions which Farley derived from a numerical solution of the kinetic equations and which were stated as items 6, 7, and 8 on pages 15 and 16, follow qualitatively from Equation (25). These conclusions do not follow from the hydrodynamic discussion.

We note, however, that Equations (25) are valid under the condition $\eta_i \ll 1$, which corresponds to the radar frequency region $f < 70$ for an altitude $h \sim 100$. Experimentally, auroral radio reflections are observed up to $f \sim 800$ MHz (which corresponds to wave numbers $\sim 3.4 \cdot 10^{-1}$, $\eta_i \sim 1$); therefore, it is necessary to generalize correctly Equations (25) to the region of large wave numbers. An attempt to generalize Farley's solutions to the region of high frequencies and large wave numbers was made in the paper [Lee, Kennel, Kindel, 1971]. As a result of the numerical solution of the dispersion equation taking into account the terms of the order of $K^2 \lambda_{De}^2$ (λ_{De} is the Debye wavelength for electrons), the authors derived the threshold nature of the dependence of the increment of high frequency waves on the electron density. It has been shown that, in the case of a fixed drift velocity of the electrons, the maximum frequency of intensifying waves increases as the electron density increases. This effect occurs at a disturbed time in the polar ionosphere when the electron density can increase by two orders of magnitude. /21

Let us summarize the main conclusions of the linear theory:

1. A threshold value of the drift velocity of electrons exists for the Buneman-Farley instabilities. This threshold value increases significantly as the aspect angle increases, and it depends on the altitude.

2. There exists an effect of the aspect sensitivity which is determinable by the dependence of the wave buildup increment on the angle ψ.

3. The dependence of the increment γ on φ, which determines the azimuthal characteristics of the radio reflections, corresponds to $\cos \varphi$, where φ is the angle subtended between the current vector and $\vec{K}$.

§ 2. Nonlinear Effects for Farley Waves

If the buildup increment of a Farley-Buneman instability is greater than zero, then the wave amplitude will increase exponentially. Since the wave amplitude, and consequently their energies, cannot increase without limit, a nonlinear mechanism restricting the amplitude growth comes into play at some time or other. /22

As is well known, a limitation to the growth of unstable waves arises due to two mechanisms — a quasilinear one and a nonlinear one. Let us briefly discuss these mechanisms separately, although generally speaking, they can operate simultaneously.

The quasilinear limiting mechanism consists of the fact that the intensifying waves so affect the particle distribution function that an effective decrease in the instability increment occurs.

In the quasilinear approximation, considering the equation for the energy of individual Fourier components of the electric field, we have the equation:

$$\frac{d|E_k|^2}{dt} = 2\gamma_k(t)|E_k|^2, \quad \gamma_k \sim A\left(U_k^2 - \frac{T_e+T_i}{M}\right). \tag{27}$$

where γ_k is the linear increment, which depends on time, because $T_e(t)$. Since the temperatures of the electrons and ions increase with time due to the interactions of the particles with the waves, the increment in the case U_k = const decreases to zero during some time interval. As a result, a certain quasi-stationary turbulence level is established, i.e., the energy density W_k of the Farley waves for a given wave number K depends weakly on the time. Estimates of the intensity of the excited waves under the quasi-stationary conditions which Kamenetskaya made (1971) give:

$$W^s = \int \frac{E_k^2}{4\pi} dk = \frac{3Q}{4\pi}\frac{B^2}{c^2},$$

where /23

$$Q = U_k^2 - 2\left(\frac{T_m}{M}\right)\left(1 + \frac{\nu_e\nu_i}{\omega_{He}^2}\frac{M}{m}\right)$$

is the degree of supercriticality which is assumed to be sufficiently small $(Q \ll U_k^2)$ that the wave spectrum is practically one-dimensional. Assuming, for example, $Q = \frac{1}{6}U_k^2$, B = 0.5 Gauss, and $U_k = 4.10^4$ cm/sec, we find $W^s \simeq 2.10^2$ ergs/sec, or $W^s/nT \sim 10^{-6}$ $(T_e \sim T_i \sim T_m$, $n \sim 3.10^5$ cm^{-3}). The energy density W_k is six orders of magnitude larger than the thermal noise energy density, and $\Delta n/n \sim 10^{-3}$, $K_{max} \sim 7.10^{-3}$ cm^{-1}, and $\Delta K \sim 10^{-3}$ cm^{-1}.

Let us emphasize once more that these estimates are valid in the case of a small supercriticality when the wave spectrum is practically one-dimensional and one can neglect the processes of nonlinear wave scattering. The characteristic time for establishing the steady state condition is $t \gtrsim 10^2$ sec. for the parameters mentioned.

The quasi-linear limiting mechanisms, which maintain the steady state condition of the Buneman-Farley instabilities, are discussed in the papers [Scadron and Weinstock, 1969]. With the use of the results of the statistical turbulent theory developed in the paper [Dupree, 1966; Weinstock, 1969; Weinstock, 1970], the kinetic equation is solved taking into account the nonlinear interaction of the waves. The solution for the ion distribution function is expressed in terms of the perturbed trajectory which describes the mean motion of ions in the turbulent field of the plasma waves.

A solution of the nonlinear dispersion equation (with the use of Farley's linear solution) gives the nonlinear decrement in Equation (2), which for the present case can be rewritten in the form

$$\frac{\partial W_{\vec{K}}}{\partial t} = \gamma_{1} W_{\vec{K}} - \Delta\omega_2 \left(W_{\vec{K}}\right) W_{\vec{K}} , \qquad (28)$$

/24

$$\Delta\omega_2 = -\frac{\nu_e \nu_i K_\perp^2 V_{Te}^2}{\omega_{Hi}^2 K^2 V_{Ti}^2} \lambda(t) ,$$

$$\lambda(t) = F\left(\sum |E_K(t)|^2\right) = \Phi\left(\int W_{\vec{K}} \, d\vec{K}\right) ,$$

$$K_\perp^2 = K_x^2 + K_y^2 ,$$

where E_K is the wave amplitude and Φ is some linear function of the turbulent energy.

As the energy of the unstable waves grows, the function $\lambda(t)$ increases, and the total cumulative increment decreases to zero upon the establishment of the quasi-stationary state. An estimate of the time to establish the quasi-stationary state gives $t = 0.45$ sec when $V_0 = 500$ m/sec.

The effect of the increasing waves on the ion distribution function is the process which maintains the quasi-stationary state. The growth of the waves corresponding to the linear buildup is compensated by their nonlinear damping by the ions.

The quasi-stationary spectrum W_k is worked out numerically in the paper [Scadron and Weinstock, 1969] for the specific value of the drift velocity $V_o = 500$ m/sec; estimates are obtained for the energy density of the waves $Wk \sim 6 \cdot 10$ ergs/cm^3, which corresponds to a maximum in the spectral energy density of the turbulence $K_{max} \sim 5 \cdot 10^{-2}$ with a spectral width in $\vec{K}$-space of $\Delta K \sim 5 \cdot 10^{-3}$. Consequently, the total energy of the turbulence under the quasi-stationary condition is, in this case, $W^s \sim W_k \Delta K \sim 10^{-12}$ ergs/cm^3, i.e., the energy of the turbulent waves is larger by four orders of magnitude than the thermal noise energy density.

A method was put forward in the paper [Rogister, 1971] which differs from the quasi-linear mechanisms discussed above for stabilizing a Farley-Buneman current instability. The stabilization mechanism in question consists of the fact that turbulence /25
having developed creates a specific flux of electrons in the current layer, which leads to the breakdown of the quasi-neutrality condition and to the development of a secondary electric field decreasing the initial drift of the electrons and the increment of the intensifying waves to zero.

We note the discrepancies in the values of the turbulence level determined by different authors. This is the value of $\Delta n/n$ derived in the paper [Scadron and Weinstock, 1969] and also in [Rogister, 1971]. Let us also note that the energy of the disturbed waves determined in the paper [Scadron and Weinstock, 1969] is larger by two orders of magnitude than it follows from the quasi-linear theory.

We note that all these estimates are significantly altered upon a variation in the parameters of the current stream. Thus, we should also point out that, although the same authors Scadron and Weinstock refer to their own theory as the theory of strong turbulence; actually, it does not resolve many difficulties in the treatment of the problems of strong turbulence and, in addition, as has been shown in Tsytovich's paper (1972), it gives incorrect results concerning the stochastic heating. An analytic solution of the nonlinear hydrodynamic equations for a somewhat idealized model of a plasma [Prasad, Sen, et al., 1971], and also a numerical solution [Sen and Prasad, 1971], has revealed the subtle effects of a current instability of the Farley type. Thus, nonlinear effects introduce wave distortions, and the wave becomes non-sinusoidal. The most rapidly growing waves with a frequency $\omega \sim 10 - 20$ kHz in the linear approximation [Kato and Hirata, 1967] have an intensity of the order of 1 — 2 v/m. A radar does not receive an individual wave ($\lambda \sim 1$ — 5 cm) directly, but a wave envelope. The energy density is larger by 1 — 2 orders of magnitude than the recorded value in the paper [Scadron and Weinstock, 1969]. /26

Since, of all the competing nonlinear stabilization mechanisms, the most significant is the one which limits the amplitude to its smallest value, one should evidently conclude that the most important of the mechanisms considered is the quasi-linear one.

However, regardless of the fact that there is a lot of controversial and contradictory material in the papers discussed, the first attempts to estimate the role of nonlinear effects in providing for the establishment of a quasi-stationary turbulence level are very successful.

Recently, comparatively weak inhomogeneities, which it is impossible to explain within the framework of Farley's current instability, since they have arisen at currents significantly smaller than the critical values necessary for the appearance of a current instability, have been detected by radar observations of the equatorial E-region [Cohen and Bowles, 1967], and also of the polar E-region [Unwin and Knox, 1972]. In connection with this fact, attempts were undertaken to explain the development of such inhomogeneities by the effect of wave interaction. A qualitative discussion of the wave interaction effect has been carried out in the paper [Dougerty and Farley, 1967]. The authors studied the occurrence of secondary waves upon the blending of two initial waves. As a result, the nonlinear interaction of the waves leads to the occurrence of waves outside the linear buildup region. Proceeding from a discussion of the conservation laws during a decay interaction,

$$\left.\begin{aligned} \vec{K} &= \vec{K}_1 + \vec{K}_2 , \\ \omega &= \omega_1 + \omega_2 , \end{aligned}\right\} \tag{29}$$

($\vec{K}_1$ and $\vec{K}_2$, ω_1 and ω_2 are the wave vectors and frequencies of the initial waves), the region of the phase velocities and possible directions of the secondary waves is determined. It is possible to explain with the help of this mechanism the motion of inhomogeneities with velocities both significantly larger and smaller than the sonic velocities. Estimates of the intensities of such secondary waves [Kamenetskaya, 1971] give a value two orders of magnitude smaller than the intensity of the initial waves. We should emphasize that the effect of radio /27 reflections on the secondary waves, the so-called four-plasma process, can occur even when the secondary waves do not actually exist (K and ω do not lie in the transparency region). The energy density of the secondary waves and the geometrical relationships were discussed in more detail on the basis of nonlinear hydrodynamics [Tsytovich, 1967] in the paper [Kamenetskaya, 1971].

§ 3. Drift Gradient Instability

The incoming fluxes of charged particles in the polar region produced additional inhomogeneities of the electron concentration. As is well known from plasma theory [Mikhailovskiy, 1971], a smooth inhomogeneity in the stable background can significantly affect the distribution and excitation conditions of the waves, particularly low frequency waves, and it results in the formation of new oscillation branches.

Here we will discuss the second type of instability, which is able to explain the formation of inhomogeneities stretched out along the magnetic field and is called the drift gradient or E and D instability, suggested for the first time in the papers [Simon, 1963; Hoh, 1963] as a model of a laboratory plasma. The presence in the ionosphere of an electric field [Mozer, 1967; Wentworch and Potter, 1970] along with a positive /28 electron density gradient in the direction of the electric field may be the cause for the origination of the unstable longitudinal waves. The possibility with respect to the ionosphere of drift gradient turbulence was treated in the papers [Reid, 1968; Tsuda and Sato, 1966; Knox, 1964, Rogister and D'Angelo, 1970; Whitehead, 1971; Whitehead, 1968; Ogawa and Sato, 1971; Knox, 1972]. It was shown in these papers that such a mechanism can be responsible for the formation of small scale inhomogeneities stretched out along the magnetic field. Just as before, we will discuss a simple hydrodynamic model for the elucidation of the main features of the instability. Let $\vec{B}$ be directed along the z-axis, and let the density gradient $\vec{\varkappa} = \mathrm{j}\,\frac{1}{n_0}\frac{\partial \ln n_0}{\partial y}$ be directed along the y-axis. Furthermore, we will discuss waves propagating along the axis ($K_x \equiv K \neq 0$). The continuity equations

$$\frac{\partial n_{1e,i}}{\partial t} + \nabla\left(n_0 \vec{V}_{1e,i} + \vec{V}_0 n_{1e}\right) = 0$$

give, in this case [see (6a)]:

$$\frac{n_{1e}}{n_0} = \frac{K V_{1ex} + i\varkappa V_{1ey}}{\omega - \vec{K}\vec{V}_{0e}}, \qquad \frac{n_{1i}}{n_0} = \frac{K V_{1ix} + i\varkappa V_{1iy}}{\omega} \qquad V_{0i} \sim 0$$

Writing out the equations of motion (3) and percentages and taking account of the fact that in an unstable plasma $\frac{\nabla n_{e,i}}{n_{1e,i}} = i\left(\frac{\nabla n_{1e,i}}{n_0} - \frac{n_{1e,i}}{n_0}\frac{\nabla n_0}{n_0}\right)$, we obtain [see Equations (9) and (10)] for the electrons:

$$\begin{aligned} 0 &= -\frac{e}{m}E_x - \omega_{He}V_{1ey} - \nu_e V_{1ex} + iKV_{Te}^2\frac{n_1}{n_0}, \\ 0 &= \omega_{He}V_{1ex} - \nu_e V_{1ey} + \varkappa V_{Te}^2\frac{n_1}{n_0} \end{aligned}$$

and, consequently, for the ions: /29

$$\begin{aligned} (i\omega + \nu_i)V_{1ix} &= \frac{e}{M}E_x + V_{Ti}^2\, iK\frac{n_1}{n_0}, \\ (i\omega + \nu_i)V_{1iy} &= \varkappa V_{Ti}^2\frac{n_1}{n_0}. \end{aligned}$$

Furthermore, substituting $n_{1e,i}$ from the continuity equation into the equations of motion, one can derive the velocity components V_{1ix}, V_{1iy}, V_{1ex}, V_{1ey} as a function of E_x. Upon substitution into Poisson equations [see Equation (14)]

$$-iKE = 4\pi e n_0\left(-\frac{KV_{1ex} + i\varkappa V_{1ey}}{\omega - \vec{K}\vec{V}_{0e}} + \frac{KV_{1ix} + i\varkappa V_{1iy}}{\omega}\right)$$

of the values of the velocities $\{V_{\perp e,i}, V_{\omega,e,i}\}$, it is easy to derive the dispersion equation being sought

$$1 = \frac{\omega_{oi}^2}{\omega(\omega - i\nu_i) - K^2 V_{Ti}^2} + \frac{\omega_{oe}^2\left(1 - \frac{\varkappa}{K}\frac{\omega_{He}}{i\nu_e}\right)}{(\omega - KV_0)\frac{\omega_{He}^2}{i\nu_e} + 2\omega_{He}\frac{K\varkappa V_{Te}^2}{i\nu_e} - K^2 V_{Te}^2} \tag{31}$$

By the method of successive approximations, we find from Equation (31)

$$\omega = \frac{\vec{K}\vec{V}_0}{1 + \frac{\nu_e \nu_i}{\omega_{He}\omega_{Hi}}}, \quad \gamma = \frac{\nu_e}{\omega_{He}^2}\frac{M}{m}\left(\omega^2 - K^2 V_S^2 + \frac{\varkappa}{K}\frac{\omega_{He}}{\nu_e}\nu_i \omega\right) \tag{32}$$

Comparing with Equations (16) and (17), we see that the expression for the frequency has not been changed, but the condition for the excitation of waves has been significantly altered and we have in place of Equation (17):

$$\left(\frac{KV_0}{1 + \frac{\nu_e \nu_i}{\omega_{He}\omega_{Hi}}}\right)^2 + \frac{\varkappa}{K}\frac{\omega_{He}}{\nu_e}\nu_i\left(\frac{KV_0}{1 + \frac{\nu_e \nu_i}{\omega_{He}\omega_{Hi}}}\right) - K^2 V_S^2 > 0 \tag{33}$$

It follows from this equation that the presence of an electron density gradient decreases the necessary critical velocity for the development of an instability. A threshold is practically absent, and the phase velocity of the waves is approximately /30 equal to the drift velocity of the electrons and can be significantly smaller than V_{Ti}.

Expressions for the increment of the unstable waves and the frequencies for the ionosphere lower region (h < 100 km) have been found in the more general case $K_x \neq 0$, $K_y \neq 0$, $K_z \neq 0$ by solution of the hydrodynamic equations using the Wentzel-Kramers-Brillouin (WKB) method from the paper [Rogister, 1970]:

$$\gamma = \frac{\omega^{(0)}}{\left(1 + \frac{\nu_e \nu_i}{\omega_{He}\omega_{Hi}}\right)}\frac{\nu_i}{\omega_{Hi}}\frac{K_x \varkappa}{K_\perp^2} + \frac{\nu_e}{\omega_{He}}\frac{(\omega^{(0)})^2 - K_\perp^2 V_S^2}{\omega_{Hi}\left(1 + \frac{\nu_e \nu_i}{\omega_{Hi}\omega_{Hi}}\right)}, \tag{34}$$

where

$$\omega^{(0)} = \frac{\vec{K}\vec{V}_0}{1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}}, \quad \varkappa = \frac{1}{n}\frac{\partial n}{\partial x}, \quad K_\perp^2 = K_x^2+K_y^2, \quad V_s^2 = \frac{T_e+T_i}{M}.$$

The first term of Equation (34) describes the drift gradient instability, which is always present at an electrode with a positive electron density gradient along the electric field. The second term of Equation (34) describes the Farley-Buneman instability, which arises at a drift velocity value of

$$V_{e0} < \frac{K_\perp}{K} V_s \left(1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}\right).$$

We note that Equation (34) agrees with Equation (32) to an accuracy of the multiplier $\left(1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}\right)$.

An investigation of the possibility of a drift gradient instability arising in the case of the polar ionosphere was carried out in the papers [Hooper, 1971; Ogawa and Sato, 1971; Knox, 1972]. A model for a polar auroral arc was considered in the form of an extended band of restricted thickness. Using the electron density distribution in the auroral arc measured by McDiarmid and McNamara (1969) and measurements of the electric field [Wescott and Stolaric, 1969], possible horizontal electron density gradients along the electric field direction were calculated. Using the results [Whitehead, 1968] for the case of an auroral arc, Knox (1972) derived the criterion for the buildup of waves propagating orthogonally to the magnetic field. The waves build up if the wave buildup increment produced by the presence of an electric field exceeds the damping due to diffusion: /31

$$\gamma = \frac{\nu_i/\omega_{Hi}}{\left(1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}\right)^2}\,\frac{\varkappa V_0}{1+\left(\frac{\varkappa}{K}\frac{\nu_i/\omega_{Hi}}{1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}}\right)^2} - \frac{K^2V_s^2\,\nu_e/\omega_{He}}{1+\frac{\nu_e\nu_i}{\omega_{He}\omega_{Hi}}} \qquad (35)$$

We note that the second term of Equation (35) corresponds to the second term in Equation (32) for the increment.

Corresponding estimates for the minimum wavelengths produced at the boundaries of the optical shapes of aurora give values $\lambda > 1$ m (for electron density gradients $\varkappa \sim 10^{-3}$ per meter, which corresponds to auroral thickness of d < 6 km). The existence of a glow thickness of some several kilometers to several meters [Maggs and Davis, 1968] and electric fields of $E \sim 10 - 30$ mv/m [Wescott and Stolaric, 1969] was assumed.

An analogous discussion of the conditions for the buildup of waves in a polar auroral arc was carried out in the papers [Ogawa and Sato, 1971; Hooper, 1971] also on the basis of the hydrodynamic equations. When recombination effects and the altitude dependence of the collision frequencies and electron density are taken into account, the estimates made by Hooper (1971) give the scale of unstable waves $\lambda > 1$ km for values of the electric fields of approximately 50 mv/m and $\varkappa \sim 5 \cdot 10^{-5}$ per meter, which corresponds to glow thicknesses less than 20 km. Estimates [Ogawa and Sato, 1971] for an analogous model of a polar auroral arc give still larger wavelength scales of $\lambda \sim$ 20 — 30 km. /32

Based on the linear theory of drift gradient instability, one can draw the following conclusions which are important for diagnostics.

1. The increments are a maximum for waves propagating parallel to the current.

2. Damping of the waves due to diffusion increases in proportion to the square of the wave numbers, $\sim K^2$ and thus, the total increment decreases sharply as K increases.

3. A threshold with respect to the current is practically absent.

4. The phase velocity of the waves can be significantly smaller than the sound velocity.

We note that the scales of the unstable waves of a drift gradient instability corresponding to the maximum increment are of the order of approximately 100 meters, which is significantly larger than is necessary for them to scatter the radar signals. Evidently, as a result of the development of turbulence, nonlinear processes can become significant in producing a transfer of the waves over the spectrum into the region of small wavelengths and, consequently, large values of K. These waves may already be responsible for auroral radio reflections. In connection with this, the investigation of possible nonlinear mechanisms which might restrict the linear growth of waves and result in a quasi-stationary spectrum is an important problem. A quasi-linear stabilization mechanism was suggested in the papers [Kim and Simon, 1969] and [Sleeper and Simon, 1970]. They /33
consider a restricted nonuniform plasma in which the insignificant excess of a certain parameter above a critical value excites a single oscillation mode of so large an amplitude that the altered electron concentration profile results in a decrease in the increment and quasi-linear stabilization.

A stabilization mechanism produced by the turbulent diffusion of ions for the drift gradient instability was investigated for application to the F region of the ionosphere in the paper [Williams and Weinstock, 1970]. The cited stabilization mechanisms were subjected to Rogister's criticism (1972), who suggested another mechanism. At the basis of this mechanism lies the nonlinear interaction of waves resulting in the transformation of the spectral energy density from the region of small wave numbers K to the region of large wave numbers K, where the waves are dense due to linear absorption.

The most complete investigation of quasi-linear effects and the effects of the interaction of unstable waves has been carried out in the paper [Sato, 1971]. It has been shown in this paper that a weakly ionized plasma in the presence of crossed fields and an electron concentration gradient may be found in two unstable states depending on the electric field. When a certain threshold value is exceeded, unstable waves build up and then are stabilized by quasi-linear effects. However, when a certain other threshold value of the parameter is exceeded, the effective interaction of modes results in another type of instability, the so-called "explosive" instability. As a result, there arises in the plasma a strong turbulence with a large number of modes interacting among themselves. The numerically determined quasi-stationary spectrum of the "explosive" /34 instability, found on the basis of the interaction of a large number of modes, is subject to a law of the form $K^{-\alpha}$ (where $\alpha \sim 2 — 4$, and K is the wave number) [Sato and Ogawa, 1972].

We will briefly enumerate the basic results of the nonlinear theory which are important for the comparison with experiments:

1. The theory predicts a suprathermal level of fluctuations in the region of wave numbers corresponding to linear stability.

2. The order of magnitude of the size of the density fluctuations for modes corresponding to the maximum of the linear buildup amounts to 25% of the basic density.

3. The quasi-stationary spectrum of drift gradient turbulence is specified by the law $K^{-\alpha}$ ($\alpha \sim 2$ — 4).

§ 4. Comparison of Experimental Data with the Models of a Radio Aurora

As has already been pointed out in the introduction, the occurrence of auroral radio reflections is caused by their scattering by inhomogeneities in the electron densities. If an electromagnetic wave $E_1 e^{i(\omega^t t - \vec{K}^t \vec{r})}$ is incident on a plasma with electron density fluctuations determined by longitudinal density waves proportional to $N_1 e^{i(\omega^\ell t - \vec{K}^\ell \vec{r})}$, then an additional current

$$\vec{J} \sim E_1 N_1 e^{i(\omega^t + \omega^\ell)t - i(\vec{K}^t + \vec{K}^\ell)\vec{r}}$$

arises in the plasma which causes a scattered electromagnetic wave. The laws of energy and momentum conservation are fulfilled:

$$\omega^t + \omega^\ell = \omega'^t, \quad \vec{K}^t + \vec{K}^\ell = \vec{K}'^t.$$

If the experimental conditions are selected in such a way that /35
the inverse scattering $\vec{K}'^t = -\vec{K}^t$ and $\omega^\ell \ll \omega^t$, then $2\vec{K}^t = -\vec{K}^\ell$, i.e., the scattering of an electromagnetic wave of wavelength λ results in a longitudinal wave with the corresponding wavelength $\lambda^\ell = \frac{1}{2}\lambda^t$, and the propagation direction is strictly opposite. In the interpretation of experimental radio reflection data, people usually compare a ring turbulence of the Farley type with type 1 inhomogeneities and with a drift gradient inhomogeneity, type 2, respectively. The separation of the various types of received radar signals becomes an important problem.

The experimental investigation of radio aurora has shown in the 1950's and 60's that the received signals are divided into two basic types, diffuse and discrete. A distinction

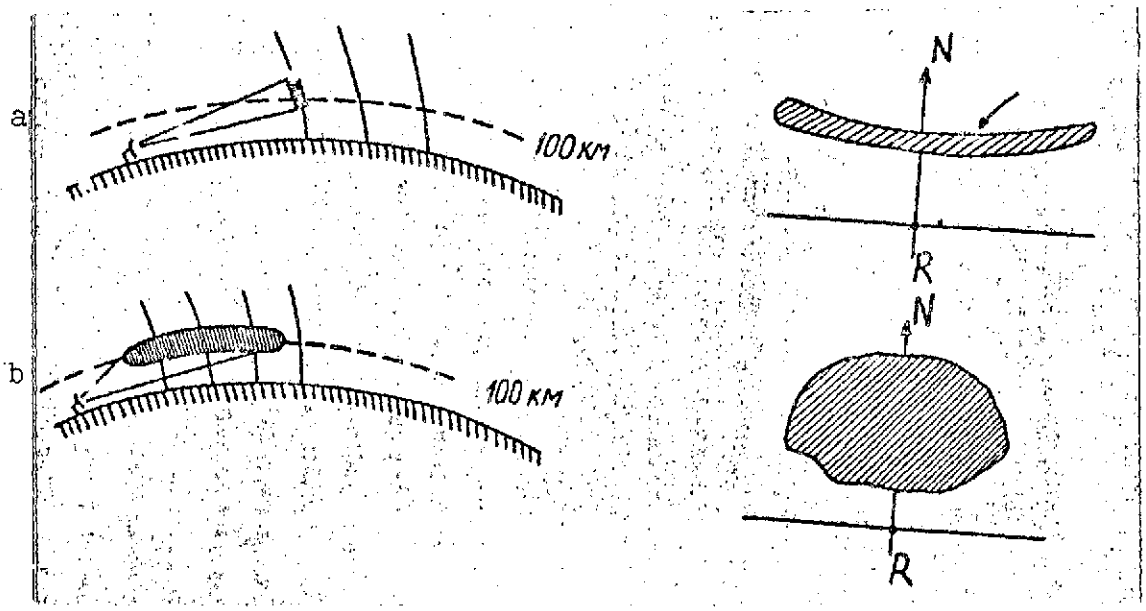

Figure 4. Schematic illustrations of the auroral reflecting regions.

a- discrete radio reflections; b- diffuse radio reflections.

either of spatio-temporal features of the reflecting regions or of the characteristics of the signals themselves was put forth as the basis of separation of the signals [Leadabrand, 1959; Presnell and Leadabrand, 1959].

The first type of reflection is of a long-lasting nature, and the reflecting region occupies a narrow altitude region and a large range interval. This type of reflection is generally called the diffuse type. The second type of reflection is of brief time duration, corresponding to a very narrow range interval and a large altitude region, and they are called discrete reflections. The qualitative distinction of the reflecting regions is shown in Figure 4.

Earlier studies of the characteristic properties of radio aurora, the aspect sensitivity, the Doppler shifts of the received signals, the azimuthal dependences, and the intensity, which was carried out by a large number of investigators, have not shown significant distinctions in these properties for discrete and diffuse radio reflections [Leadabrand, 1962; Presnell and Leadabrand, 1959; Collins, 1951; Leadabrand, 1959; Unwin, 1959]. /37
However, a difference of the diurnal paths of visibility of diffuse and discrete reflections was clearly revealed. Thus, the maximum probability of appearance of discrete radio reflections occurs at times near midnight [Forsyth, 1968; Unwin and Knox, 1968], while the diffuse radio reflections occur in the evening and morning hours [Unwin, 1959; Leadabrand, 1962]. Unwin (1959) divided up the radio reflections into four types: diffuse, diffuse with structure, brief discrete, and lengthy discrete.

At the present time, in agreement with the IAGA recommendations [IAGA, 1968; Sverdlov, 1971], it has been proposed to divide the signals into type B_1, diffuse and diffuse with structure, and type B_2B_3, discrete reflections differing in duration: $B_2 < 1$ min and $B_3 > 1$ min. In agreement with this division, people are presently attempting to compare the different types of radio reflections and different mechanisms for the formation of inhomogeneities stretched out along the Earth's magnetic field. Thus, the diffuse type of reflections, type B_1, is compared with inhomogeneities of type I, and the discrete type of reflections, type B_2B_3, is compared with type II inhomogeneities [Unwin and Knox, 1971; Knox, 1972; Unwin and Baggaley, 1972]. We will list how the properties of both types of signals, which arise due to a difference in the properties of the inhomogeneities formed as a result of the development of the two types of instabilities, differ:

1. Inhomogeneities arising as the result of the development of a Farley instability are localized in a narrow altitude region [Unwin and Knox, 1968; McDiarmid, 1970], and the inhomogeneities caused by the drift gradient instability occur in a wider altitude region [Unwin and Baggaley, 1972; Knox, 1972]. The experimental confirmation of the correspondence of diffuse reflections to the Farley instability was provided by Gadsden (1967) and Leadabrand and Schlobohm (1968). /38

The altitude and thickness of the reflecting layers for the discrete reflections B_2B_3 correspond according to measurements [Unwin, 1959; Presnell and Leadabrand, 1959] to the region of maximum increments of the drift gradient instability [Knox, 1972].

2. There occurs for type B_1 signals a threshold dependence of the appearance probability on the magnetic perturbing ability of the existence of a critical current, but there is no threshold dependence for type B_2B_3 signals. The correlation of B_1 signals with ring systems has been confirmed in numerous papers (Ecklund and Unwin, 1971; Gadsden, 1967; Lange-Hesse, 1968; Harris and Kavadas, 1970; and others]. For example, the explanation of the diurnal variation in the appearance of diffuse reflections [Unwin and Knox, 1968; Unwin, 1966] has been given in terms of a Farley instability with the use of a specific ring model.

The presence of an electron density gradient in the direction orthogonal to the magnetic field is essential in order for a drift gradient instability to arise; therefore, a correlation of this type of radio reflection with the incoming streams of particles capable of creating such gradients should be observed. Investigations of the fine structure of the particle dumping have determined, according to the observations [Maggs, 1968], the sizes of the horizontal gradients sufficient to explain the development of type B_2B_3 radio reflections due to a drift gradient instability [Knox, 1972]. A correlation of the radio reflection with the structure of the electron dumping was detected experimentally by Leadabrand [Leadabrand and Hodges, 1967] at a frequency of 139 MHz; however, they did not detect such a correlation with the proton intrusions. In contrast to this result, a spatial coincidence of the regions of proton dumping and the radio reflections has been noted in the paper [Hagfors and Jonson, 1971] upon a comparison of the satellite measurements of electron and proton dumping with the radio reflections. However, these measurements are of a random nature; moreover, no separation of the signals was carried out. /39

The quasi-periodic fading of the 42 MHz signals was detected in measurements described in the paper [Sofko and Kavadas, 1969], in which it is confirmed that the radio reflections come about from limited layers of electron density produced directly by the electrons pouring in. It was remarked that the periodicity arises from the periodicity of the primary stream, and Unwin (1971) confirmed that these observations can be explained on the basis of the drift gradient instability.

The coincidence of the region in which radio reflections are observed with the auroral oval [Unwin, 1966; Shipstone, 1969], where the dumping of particles maintains the necessary horizontal gradients, can also be qualitatively explained by the development of a drift gradient instability in these regions.

3. The buildup increment of the two types of instabilities depend in different ways on the wavelengths. There is a weak wavelength dependence for the Farley instabilities, and the increment sharply decreases as K increases for the drift

gradient instability and therefore, on the basis of this determination of the theory, one should expect a larger amount of type B_2B_3 signals at the lower frequencies. Similar estimates were made in Unwin's papers [Unwin and Knox, 1971], in which it was shown that there exists a tendency for the ratio of the amount of B_2B_3 signals to B_1 to decrease as the wavelength decreases.

4. The Doppler shift of the type B_1 signals should, on the basis of the Farley instability, correspond to a sound velocity in the plasma on the order of V_{Ti}, i.e., be somewhat larger than the thermal velocity of the ions. The Doppler shift for B_2B_3 signals can be significantly smaller V_{Ti}. However, the experimental determinations of the size of the Doppler shift of the received signals are of a contradictory nature. /40

The measurement of the Doppler shifts noted in the papers [Leadabrand and Schlobohm, 1965; Hofstee and Forsyth, 1969] has corresponded completely to the Farley instability model. Within the framework of a radio aurora model [McDiarmid and McNamara, 1969; McDiarmid, 1970], the spectral characteristics of the signals received by Hofstee [Hofstee and Forsyth, 1969] are completely explained upon the assumption of the existence of an electric field of the order of 30 mv/m in the polar ionosphere. The distribution of the Doppler shifts at a frequency of 1295 MHz, which have a maximum near 4 kHz corresponding to a sound velocity on the order of 400 m/sec, has been derived by Abel (1969). However, along with the numerous confirmations of the presence of signals corresponding to the Farley instability model, a large number of signals have been detected with a Doppler shift corresponding to speeds both significantly larger and significantly smaller than the sound velocity in the plasma.

The difficulties of separating the various types of reflections with respect to their Doppler shifts arise because of the indeterminancy in the interpretation of the shifts corresponding to speeds significantly smaller than the sound velocity. On the other hand, small Doppler shifts of the signals can be explained in terms of ionosonic secondary waves [Dougherty and Farley, 1967], or, on the other hand, as a mixture of primary and secondary waves within the framework of drift gradient turbulence. /41

One should note that the explanation of the Doppler shifts of type B_2B_3 signals as a mixture of primary and secondary waves encounters difficulties at small drift velocities, when the excitation criterion of the Farley instability is not satisfied. A comparison of type B_2B_3 signals to the inhomogeneities arising as a result of the development of a drift gradient instability is very promising.

Unwin's study [Unwin and Knox, 1968] of the dependence of the velocity of the inhomogeneity motion on local time shows agreement of B_1 signals of an ionosonic Farley instability and a clear disagreement of type B_2B_3 signals, which correspond to inhomogeneities moving with the electron drift velocity.

We note that a comparison of the two types of inhomogeneities and types I and II radio reflections encounters a number of difficulties. These difficulties arise both due to the insufficiently developed nonlinear theory for the formation of inhomogeneities and also because of the complexity of separating the variety of signals by types. Such a separation of the signals has now been satisfactorily carried out with respect to the equatorial region [Balsley, 1970; Balsley and Farley, 1971]. However, the situation in the polar ionosphere is far more complex; even the traditional separation of the signals into diffuse and discrete types is difficult. This situation is associated both with the variety of types of radio aurora signals and with the complexity of interpreting the experimental data in view of the complicated geometry of radio aurora [Uspenskiy, 1972; Pyatsi and Sverdlov, 1971; Ponomarev and Vershinin, 1967; Bagaryatskiy, 1961].

We also point out that the theory of the formation of type I inhomogeneity has been worked out with respect to the equatorial region and has only been adapted to the conditions in the polar ionosphere. Thus, Moorcroft (1972a, b) has attempted to explain, on the basis of the actual conditions of an auroral electrode, the basic nature of a radio aurora — the possibility of large aspect angles. An attempt to take into account the vertical electron density gradients was unsuccessful. However, the conclusion of possible large drift velocities (up to 5 km/sec) has permitted explaining Hofstee's results. /42

An attempt was made in the paper [Volosevich and Liperovskiy, 1973] to take into account the ionosphere turbulence due to the presence of longitudinal currents exciting ionosonic turbulence. The taking into account of nonlinear scattering of Farley waves by ionosonic waves results in a change in the dependence of the buildup increment of the waves on the aspect angle. Due to this effect, it becomes possible to explain the cases of aspect angles $\psi \sim 10°$ at times of the ionosphere agitation.

The theory for the origin of type II inhomogeneities has, on the contrary, been developed with application to the radio aurora. Thus, on the basis of the expression for the drift gradient instability increment, one can determine the critical minimum thickness of the arc specifying the threshold for the origin of radio reflections from type II inhomogeneities [Volosevich and Liperovskiy, 1972]. However, there remains in this direction an open problem as to the production of waves with large values of K. One can resolve this problem only by taking into account the different nonlinear processes and with the production of a quasi-stationary spectrum of the drift gradient turbulence.

Regardless of the indicated contradictions and difficulties in comparing the different kinds of inhomogeneities and types /43
of signals of radio aurora, the investigation of the different physical processes constituting the basis of the formation of inhomogeneities should stimulate the division of radio reflections into types. Also, the refinement of the methods of radar probing of the polar ionosphere and the possibility of a clear separation of signals will permit the growth of theoretical investigations. In this way, the intensively developed method of parametric excitation of normal waves in the ionosphere is a very promising one at the present time. [Lee, Kaw, and Kennel, 1972].

Conclusions

The analysis which has been carried out shows that the /44
investigation of all possible types of instabilities in the polar aurora associated with the presence of strong currents, streams of incident particles, and electron density gradients has made it possible to construct a still crude model of a radio aurora which explains the properties of auroral radio reflections. On the other hand, the appearance in the course of experiments of different types of radio reflections and their properties, which correspond to different types of instabilities, results in a refinement of the theoretical models.

Besides the cited types of instabilities capable of causing auroral radio reflections, it is very probable that there exist still other non-investigated instabilities associated, for example, with the deformation of the ion and electron distribution functions due to the incoming streams, with currents parallel to the magnetic field, with spatial variations in the magnetic field, temperature, and electron density, and with the variation of the drift velocity without these. Upon comparing the experimental data with the radio aurora model, significant difficulties occur which arise both due to the insufficiently developed nonlinear theory of the turbulence of Farley waves and of the drift gradient instability and because of the inaccuracy in signal classification. Regardless of this, the further study of instabilities well known in plasma physics is applicable to the conditions of the polar ionosphere. The segregation of the different types of radio reflections and a comparison of them with the properties of turbulence arising as a result of the development of this or the other types of instability is necessary for the creation of a more satisfactory model of the radio aurora.

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